Intermediate Scale Inflation and Metastable Supersymmetry Breaking

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Abstract

We investigate the possibility of obtaining a low scale of supersymmetry breaking within the ISS framework using a metastable vacuum. This is achieved by introducing an R symmetry preserving gravitational coupling of the ISS sector to a relatively low scale inflationary sector. We find the allowed range for the supersymmetry breaking scale, 10^4 GeV $\lesssim \mu \lesssim 10^8$ GeV, which is low enough to be amenable to gauge supersymmetry breaking mediation. This scenario is based upon a so-called hilltop inflation phase whose initial condition problem is also addressed.

It has been recently realised by ISS [1] that supersymmetry breaking can be achieved in a metastable vacuum which is separated from the true supersymmetry preserving vacuum by a barrier that can guarantee a life-time for the false vacuum which exceeds the age of the universe. One particular advantage of this setting is that the IR free magnetic description (which is dual to a UV free electric theory) is suitable to study low energy physics. This opens up the possibility of describing supersymmetry breaking at low energy (the ISS scale μ) compared to the Landau pole of the magnetic phase. Within the metastable supersymmetry breaking framework, it has been recently shown in [2] that R symmetric gravitational couplings between the supersymmetry breaking sector and the inflation one would help determining the ISS scale. It would also provide a natural explanation for why the universe should end up in the metastable minimum instead of the supersymmetric minimum.

The connection between supersymmetry breaking and inflation may shed some light on our understanding of scales beyond the standard model of particle physics. Indeed cosmological observations of the cosmic microwave background anisotropies single out a very large scale close to the GUT scale when interpreted within the inflationary paradigm. The magnitude of the temperature fluctuations is given by the height of the Sachs-Wolfe plateau and corresponds to $\frac{\delta T}{T} \simeq 6.6 \times 10^{-6}$ [3, 4]. This translates into a constraint on the inflationary potential V_I,

$$\left(\frac{V_I}{\epsilon}\right)^{1/4} \simeq 6.6 \times 10^{16} \text{ GeV}, \tag{1}$$

where ϵ is the slow roll parameter defined as $\epsilon = \frac{M_P^2}{2}\left(\frac{V'}{V}\right)^2$. So depending upon the value of ϵ at the time of horizon exit, the inflationary scale ($V_I^{1/4}$) can be estimated. For example, in supersymmetric hybrid inflation [5] ($\epsilon \sim 10^{-8}$), this scale turns out to be 10^{15-16} GeV, *i.e.* $\sim$ GUT scale. We have found in [2] that this corresponds to the scale of supersymmetry breaking in the range of gravity mediation. For sufficiently low ϵ, this characteristic scale would be lower. In particular, we will find that for intermediate values of $V_I^{1/4} \approx 10^{11}$ GeV, the supersymmetry breaking scale could be as low as 10^4 GeV. Of course, this is within the right ball park for gauge mediation of the supersymmetry breaking to the MSSM. This could well be hint that supersymmetry breaking effects appear at low energy and could be observable at the LHC.

In this letter our aim is to find a metastable supersymmetry breaking at a low scale (μ, the ISS scale) which is consistent with gauge mediation of supersymmetry breaking. Following the approach in [2], we assume the existence of two sectors, the inflation and the ISS sector, which communicate with each

other only through gravity[1] (respecting $U(1)_R$ symmetry). The inflation sector consists of superfields $(\chi, \bar\chi)$ and S, with R charges, $R[\chi\bar\chi] = 0$ and $R[S] = 2$. As we are dealing with gravitational interactions between the two sectors specified above, it is quite natural to consider inflation models in the framework of supergravity. Therefore we must specify the Kähler potential of the inflation sector. We assume that the Kähler potential is invariant under a shift symmetry of the inflaton chiral multiplet $\chi, \bar\chi$. The choice of this shift symmetry is mainly motivated by the solution to the η problem[2]. Thus the inflaton direction $\chi = \bar\chi$ does not receive any mass-squared term $\sim O(H_I)$ the presence of which otherwise would spoil the flatness of the potential. The Kähler potential is given by

$$K_- = |S|^2 + \frac{1}{2}|\chi - \bar\chi|^2 + a_1\frac{|S|^4}{4M_P^2} + a_2\frac{|\chi - \bar\chi|^4}{4M_P^2} + a_3|\chi - \bar\chi|^2\frac{|S|^2}{2M_P^2}, \tag{2}$$

where we keep higher order terms whose necessity will be spelt out later.

It is a known fact that to generate inflation we need to break the exact shift symmetry in order to give a slope to the inflaton potential. This is achieved by introducing a higher order (*gravitational*) term in the superpotential[3]. Now, the inflaton field is defined by $\chi = \bar\chi$, while the $\chi = -\bar\chi$ direction corresponds to a massive field which plays no role in inflation and can be discarded from the discussion. Hence we keep only the inflaton field which, for convenience, we still denote χ, in the inflationary superpotential that we write in the form

$$W_{inf} \quad = \quad S\left(k\frac{\chi^n}{M_{\rm P}^{n-2}} - M^2\right), \tag{3}$$

where $n > 2$ and a discrete symmetry identically transforming χ and $\bar\chi$ guarantees the form of the superpotential.

The ISS sector is described by a supersymmetric $SU(N_c)$ gauge symmetry with N_f flavors of *massless* quark-antiquark pairs in the electric theory. Here Λ is the strong-coupling scale of the theory, below which the theory can be described as the magnetic dual, $SU(N)$ gauge theory, where $N = N_f - N_c$ with N_f flavors of magnetic quarks, $q_i^c, \tilde q_c^i$, ($i = 1\ldots N_f$ and $c = 1\ldots N$) and a $N_f \times N_f$ gauge singlet superfield Φ_j^i (the meson field $\Phi = Q\tilde Q/\Lambda$). The magnetic theory is IR free if $N_c + 1 \leq N_f \leq \frac{3}{2}N_c$ and has a superpotential given by

$$W = h\mathrm{Tr}q\Phi\tilde q, \tag{4}$$

for massless quarks, along with the dynamical superpotential

$$W_{dyn} = N\left(h^{N_f}\frac{\det\Phi}{\Lambda^{N_f-3N}}\right)^{\frac{1}{N}}, \tag{5}$$

where $h = O(1)$. The R-charges are such that Φ has a R-symmetry charge $R[\Phi] = 2$, $R[Q] = R[\tilde Q] = 1$ up to a baryon number and $R[q, \tilde q] = 0$.

The interaction between these two sectors can be described (in the magnetic phase) by

$$W_{int} = \lambda\frac{\chi^n}{M_P^{n-1}}\Lambda\mathrm{Tr}\Phi \tag{6}$$

which respects $U(1)_R$ symmetry as well as the discrete symmetry imposed upon the $\chi, \bar\chi$ fields (this restricts also the form of the χ-dependent terms in the superpotential of the inflation sector as we discussed before.). Once inflation ends, the χ field gets a vev and W_{int} induces a mass term for the electric quarks,

$$W_{\rm ISS} = h\mathrm{Tr}q\Phi\tilde q - \mu^2\mathrm{Tr}\Phi, \tag{7}$$

which is the same superpotential as analysed by ISS with μ^2 defined as $\frac{\lambda}{k}\frac{\langle\chi^n\rangle}{M_P^{n-1}}\Lambda$. It turns out that for $\mu \ll \Lambda$, supersymmetry is broken at the metastable minimum, $\langle\Phi\rangle = 0$, $\langle q\rangle = \langle\tilde q\rangle = \mu$. In our approach,

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[1]How the ISS sector interacts with the MSSM, so that the soft supersymmetry breaking effect can be seen, is beyond the scope of our present work. For recent works in this direction, see [6, 7, 8].

[2]The shift symmetry is also an essential ingredient of some string inflation models such as the ones based on the compactification manifold $K_3 \times T^2/Z_2$ where the free motion of branes along the two torus is translated as a shift symmetry in the Kähler potential[9].

[3]A higher order breaking of the shift symmetry will also be present in the Kähler potential and will be crucial in getting rid of the initial condition problem for hilltop inflation.

the scale of supersymmetry breaking can be written as

$$\mu^2 = \sqrt{3}\frac{\lambda}{k}H_I\Lambda, \qquad (8)$$

where H_I is the Hubble scale during inflation ($H_I^2 = \frac{V_I}{3M_P^2}$) as the χ^n term in the inflationary superpotential cancels the vacuum energy during inflation (M^4). We have assumed that gravity respects the R-symmetry as well as the discrete symmetry imposed upon χ and $\bar{\chi}$. Using the constraint on μ from the metastability condition [1] ($\mu < \Lambda$), we find that $H_I < \Lambda$. This means that the only way of achieving a low scale of supersymmetry breaking is through low value of H_I, i.e. lowering the scale of inflation. With supersymmetric hybrid inflation model, it is not possible to lower the scale of inflation very much [2]. On the contrary if we adopt a hilltop type of inflation model where the inflaton rolls down from a saddle point towards a minimum, we can achieve a low value of H_I. This leads to a low value of V_I which is then consistent with the COBE data [3] as ϵ turns out to be very small (see Eq.(1)). This model has also the power of explaining a low value of the spectral index as obtained from the WMAP 5 years data [4] $n_s \approx 0.96 \pm 0.014$. We find that a minimalistic choice for n is 4. This entails that the discrete symmetry we have discussed before would be a Z_4 invariance, under which both χ (also $\bar{\chi}$) carry charges i while S has charge 1.

We are now going to discuss the inflationary scenario in more detail. We start with the superpotential in Eq. (3) for $n = 4$. The supersymmetric vacuum is given by $S = 0$ and $k\chi^4 = M^2M_P^2$. The scalar potential in supergravity using the Kähler potential in eq.(2) along the inflaton direction reads

$$V \simeq M^4\left[1 - a_1\frac{|S|^2}{M_P^2} - 2k\frac{|\chi|^4}{M_X^4} + k^2\frac{|\chi|^8}{M_X^8}\right], \qquad (9)$$

where $M_X^2 = MM_P$ and we have chosen $a_1 < -1/3$, so that S receives a positive mass square greater than $H^2 \simeq M^4/3M_P^2$ during inflation and therefore rapidly settles to zero. Such a class of potentials[10] has been already considered [11] and happens to be a good approximation to the dynamics of racetrack inflation in string theory[12]. Inflation takes place when the field starts close to the origin ($\chi \simeq 0$) where the potential is maximal. From there it rolls down at a slow rate before eventually settling down at the supersymmetric minimum far away from the origin, $k\langle\chi^4\rangle = M_X^4$. The fact that the inflaton starts from a low value compared to the Planck scale is an initial condition issue which will be discussed later.

The slow roll parameters are given by (for $|\chi| \ll M_X$)

$$\epsilon = \frac{M_P^2}{2}\left(\frac{V'(\chi)}{V(\chi)}\right)^2 \simeq 32M_P^2\frac{\chi^6}{M_X^8}, \qquad (10)$$

$$\eta = M_P^2|\frac{V''(\chi)}{V(\chi)}| \simeq 24M_P^2\frac{|\chi|^2}{M_X^4}. \qquad (11)$$

The field value at the end of inflation, χ_f, is given by $|\eta| \simeq 1$,

$$\chi_f \simeq \frac{1}{2\sqrt{6k}}\frac{M_X^2}{M_P}. \qquad (12)$$

The number of e-foldings, N, then relates the initial value of the inflaton field, χ_0 at the time of horizon exit with χ_f by

$$N = \frac{1}{M_P^2}\int_{\chi_f}^{\chi_0}\frac{Vd\chi}{V'} \simeq \left(\frac{1}{\chi_0^2} - \frac{1}{\chi_f^2}\right)\frac{M_X^4}{16kM_P^2}, \qquad (13)$$

hence $\chi_0 \simeq \frac{1}{2\sqrt{2k}(3+2N)^{1/2}}\frac{M_X^2}{M_P}$ where we have used eq.(12). The spectral index is given by

$$n_s \simeq 1 - 2\eta \simeq 1 - \frac{6}{3 + 2N}. \qquad (14)$$

With $N = 52$, the resulting[4] spectral index $n_s \simeq 0.945$ which is within 1σ of the central value of the spectral index as recently prescribed by the WMAP result [4].

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[4]The number of e-foldings is related with the scale of inflation by $N \simeq 60 - \log\left(\frac{10^{16}\,\text{GeV}}{V_I^{1/4}}\right)$.

The inflation scale is determined by the COBE normalisation

$$(\frac{V_I}{\epsilon})^{1/4} \sim 6.6 \times 10^{16} \text{ GeV}. \tag{15}$$

Using eqs.(9) and (10), ϵ at the time of horizon exit can be expressed as

$$\epsilon \simeq \frac{k^2}{16 \left(3 + 2N\right)^3} \left(\frac{M_X}{M_P}\right)^4. \tag{16}$$

Therefore using $V_I = M^4$ and $M_X^2 = MM_P$, we find from eq.(15) that the scale of inflation is

$$M \simeq \frac{4.5}{(3 + 2N)^{3/2}} \times 10^{14} \text{ GeV} \sim 10^{11} \text{ GeV}. \tag{17}$$

Notice that the natural scale $\approx 10^{14}$ GeV is reduced, thanks to the e-fold factor $(3 + 2N)^{3/2} \approx 10^3$. The initial field value χ_0 is required to be $\chi_0 \sim 10^{-4} M_X$, this initial condition issue will be discussed at the end of this letter.

Once inflation is over, the coupling in eq.(6) implies that the SUSY breaking scale $(F_\Phi = \mu^2)$ is given by

$$\mu^2 = \frac{\lambda}{k} \frac{M_X^4 \Lambda}{M_P^3}. \tag{18}$$

In terms of the Hubble rate during inflation, $H_I \simeq \frac{M^2}{\sqrt{3}M_P}$ (in the inflationary scenario considered above, $H_I \sim 10^4$ GeV), this leads to eq.(8). In order to maintain the metastability condition in the ISS sector, one has to impose a constraint $\mu < \Lambda$, which in turn sets a lower bound (along with eq.(18)) on the scale of supersymmetry breaking as

$$\mu > H_I \simeq 10^4 \text{GeV}, \tag{19}$$

for $\lambda/k \sim 0(1)$. In the following we will obtain an upper bound while discussing reheating at the end of inflation. In a similar fashion to [2], Φ is also stuck at origin during inflation due to the presence of a mass term bigger than H_I due to the supergravity corrections. Notice that when inflation is over, this point, $\Phi = 0$, becomes a local minimum (this supersymmetry breaking minimum appears when μ becomes non-zero as a result of displacement of χ from the top of the potential in the inflation sector) and so the field does not move. This explains why the universe should prefer the supersymmetry breaking minimum rather than the supersymmetric one in the ISS sector when one considers the evolution of the universe.

At the end of inflation, the inflaton field performs damped oscillations about the supersymmetric minimum of the inflation sector and decays. The main decay channel follows from

$$V \supset |\frac{\partial W}{\partial \Phi}|^2 = |hq\tilde{q} + \lambda \frac{\chi^4}{M_P^3} \Lambda|^2. \tag{20}$$

This leads to the decay of χ into magnetic quarks (since we are already in the magnetic phase) with the decay width

$$\Gamma \simeq \frac{h^2 \lambda^2}{8\pi k^{3/2}} \frac{\Lambda^2}{m_\chi} \left(\frac{M_X}{M_P}\right)^6, \tag{21}$$

where m_χ is the mass of the inflaton, $m_\chi = \sqrt{2k}M^2/M_P$. Thus the reheat temperature $T_R \simeq \frac{1}{7}\sqrt{\Gamma M_P}$ is given by

$$T_R \simeq \frac{h\lambda\Lambda}{14\sqrt{2\sqrt{2}\pi}} \frac{1}{k^{7/8}} \left(\frac{M}{M_P}\right)^{3/4}. \tag{22}$$

Imposing that reheating should take place before the electroweak transition[5] $T_R \gtrsim 10^2$ GeV leads to a lower bound $\Lambda \gtrsim 10^8$ GeV where we have used eq.(17) and $h \sim \lambda = O(1)$. Since from the metastability condition we know $\mu < \Lambda$, it results into an upper bound on the SUSY breaking scale, $\mu < 10^8$ GeV,

[5]We impose this condition keeping in mind that the most popular way of generating the baryon asymmetry of the universe via leptogenesis will be ineffective if the reheat temperature after inflation is less than 10^2 GeV. This particular decay chanel we consider does not produce MSSM particles and for that the inflaton sector should couple with the MSSM sector which we will not discuss here.

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obtained for the lowest value of Λ. Combining it with eq.(19), we find that our scenario constrains the scale of supersymmetry breaking as follows

$$10^4 \text{ GeV} \lesssim \mu \lesssim 10^8 \text{ GeV}. \tag{23}$$

In this work, we have not focused on the mediation mechanism, *i.e.* how the supersymmetry breaking will be mediated to the MSSM sector. We keep this for future work where we will deal with inflation and a deformed ISS model of supersymmetry breaking in order to include R symmetry breaking also.

Let us now come back to the initial condition problem mentioned previously. Indeed we have assumed that χ is small initially, $\sim 10^{-4} M_X$. This calls for an explanation. A first possibility springs from the fact that prior to inflation, the universe could be radiation dominated and in a high temperature phase. Here we present a mechanism following [13] which leads to a satisfactory explanation for the initial condition problem. To address the initial χ value, we introduce one or more superfields Y_i with $R[Y_i] \neq 0$. They may have interactions with the MSSM (or extended MSSM) superfields. We also postulate a higher order shift symmetry breaking term in the Kähler potential which is actually a cross term between[6] Y_i and $(\chi + \bar{\chi})$,

$$K_+ = \sum_{i=1}^{P} b_i \frac{|\chi + \bar{\chi}|^2}{2M_{\text{P}}^2} |Y_i|^2. \tag{24}$$

Following the approach in [13], the above term leads to an interaction, the thermal average of which is given by

$$b_i \langle \partial_\mu Y_i \partial^\mu Y_i \rangle \frac{\chi^2}{M_P^2} \simeq b_i m_{Y_i}^2(T) \frac{T^2}{12} \frac{\chi^2}{M_P^2}, \tag{25}$$

where $m_{Y_i}^2(T)$ is the thermal mass for the Y_i field which depends on all the other interactions of Y_i. For instance a coupling to matter fields f and $\bar{f}$ in a Yukawa-like fashion $W \supset \gamma_i Y_i f \bar{f}$ leads to a thermal mass $m_{Y_i}^2(T) = \frac{\gamma_i^2}{6} T^2$. This is larger than the Hubble rate ($H \sim \frac{T^2}{M_P}$ in the radiation dominated era) and drives Y_i to the origin. As a result, the inflaton χ gets an effective mass square, $m_{\text{eff}}^2 \simeq \sum_{i=1}^{P} \sigma_i T^4/(12M_P^2)$, where for instance $\sigma_i = b_i \gamma_i^2/6$, which is related to the Hubble mass squared as $m_{\text{eff}}^2 = p^2 H^2$ in the radiation dominated pre-inflationary epoch. Therefore solving the evolution equation for χ, one finds[7]

$$\chi = \chi_* \left(\frac{\rho}{\rho_*} \right)^{1/2} \cos \left(\sqrt{p^2 - 1/4} \ln \left(\frac{R}{R_*} \right) \right), \tag{26}$$

where χ_*, R_* and ρ_* represent the amplitude (supposed to be $\sim M_X$), the scale factor and the energy density when the Y_i fields are in thermal equilibrium at temperature T_*. The above expression implies that once $T < T_*$, the χ field performs damped oscillation about $\chi = 0$ and it would continue till the vacuum energy of inflation is comparable to the radiation density, i.e. $\rho \sim M^4$. If the amplitude at this point coincides with the initial value required for χ, $\chi_0 \sim 10^{-4} M_X$, then the initial condition issue is resolved and inflation starts. Equating $\chi_0 = c M_X \left(\frac{M}{T_*} \right)^{1/2}$ where $c = \langle \chi_* \rangle \cos \left(\sqrt{p^2 - 1/4} \ln \left(\frac{R_0}{R_*} \right) \right) \lesssim 1$ includes the ambiguity of the field value of χ_* as well as the value of the oscillatory cosine term, it follows that with $T_* \sim 10^{18-19}$ GeV the initial value problem is solved. Hence for an initial temperature close to the Planck scale, the initial value of χ is such that the field starts rolling slowly and leads inflation at the end of pre-inflationary phase. Note that the inflaton gets a purely thermal mass through the breaking of shift symmetry before inflation, such a term becoming negligible during inflation so that the shift symmetry conserving term still protects the inflation mass and solves the η problem[8].

It is worth comparing the present set up with the one in [2]. Both are based on R-symmetry, and the inflationary superpotential is basically the same, but thanks to the differences in their Kähler potentials, the inflaton is S in [2] and $(\chi + \bar{\chi})$ in the present work, and the inflation mechanisms are different. As a consequence, the inflationary models in [2] yield a supersymmetry breaking scale consistent with gravity

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[6]The shift symmetry preserving term, although present, will not intervene as the inflaton direction is $\chi = \bar{\chi}$.

[7]Other fields may also have a thermal mass, but those are irrelevant for our analysis as they are not destabilising anything.

[8]During the pre-inflationary era, the Y_i fields are driven to the origin as their thermal masses are larger than the Hubble rate. At the end of this pre-inflationary epoch and as soon as inflation starts, the Y_i fields have a mass term of order H_I which guarantees their stability at the origin. As a result, the non-renormalisable term K_+ does not contribute to the η problem despite its shift-symmetry breaking feature.

mediation, while here this scale is much lower and consistent with gauge mediation. Although we do not tackle the issue in this paper, the needed couplings of the inflation and supersymmetry breaking to the MSSM fields are expected to be different as well. We have studied the gravitational coupling between intermediate (or even low) scale inflation and ISS metastable supersymmetry breaking in models where both phenomena are regulated by an R symmetry. We have found how the supersymmetry breaking scale is related to the low value of the Hubble rate during inflation. Requiring the reheating temperature to be above the electroweak scale, we obtain bounds on the supersymmetry breaking scale as 10^4 GeV $\lesssim \mu \lesssim 10^8$ GeV.

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